

Exotic baryons from a heavy meson and a nucleon

– Positive parity states –

Yasuhiro Yamaguchi¹, Shunsuke Ohkoda¹, Shigehiro Yasui², and Atsushi Hosaka¹

¹*Research Center for Nuclear Physics (RCNP),*

Osaka University, Ibaraki, Osaka, 567-0047, Japan and

²*KEK Theory Center, Institute of Particle and Nuclear Studies,*

High Energy Accelerator Research Organization,

1-1, Oho, Ibaraki, 305-0801, Japan

Abstract

We study heavy baryons with exotic flavor quantum numbers formed by a heavy meson and a nucleon ($\bar{D}N$ and BN) with positive parity. One pion exchange interaction, providing a tensor force, dominates as a long range force to bind the $\bar{D}N$ and BN systems. In the heavy quark mass limit, pseudoscalar meson and vector meson are degenerate and the binding mechanism by the tensor force analogous to that in the nuclear systems becomes important. As a result, we obtain the $\bar{D}N$ and BN resonant states in the $J^P = 1/2^+$, $3/2^+$ and $5/2^+$ channels with $I = 0$.

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I. INTRODUCTION

The recent finding of the twin Z_b resonances near the $B\bar{B}^*$ and $B^*\bar{B}^*$ thresholds [1–3] has added a new evidence of the exotic states in addition to candidates such as f_0 , a_0 and $\Lambda(1405)$ in strangeness sector [4–6], X , Y and Z in charm and bottom sectors [7–9], implying hadronic composites or molecules. The appearance of the states near the threshold is a necessary condition that the states can be interpreted as hadronic composites with keeping their identities as constituent hadrons. The mechanism of forming hadronic composites near the threshold depends on the nature of the interaction among the constituent hadrons.

In this respect, the one pion exchange potential is of great interest as one of the most important meson-exchange potentials between the two constituent hadrons [10–13]. The one pion exchange naturally works when the constituent hadrons have non-zero isospin values. We note that the existence of the pion is a robust consequence of spontaneous breaking of chiral symmetry [14]. A unique feature of the one pion exchange potential is the tensor force due to the pseudoscalar nature of the pion. The tensor force mixes the states with different angular momentum, i.e. L and $L \pm 2$. This causes a mixing of the different configurations in a hadronic state and thus yields an attraction between the two constituent hadrons with lower L state. In fact, it is known that the tensor force is the leading mechanism of the binding of atomic nuclei [15].

The pion exchange is possible also for the heavy hadron systems containing heavy pseudoscalar meson $P = \bar{D}$, B and heavy vector meson $P^* = \bar{D}^*$, B^* . The Yukawa vertices of $PP^*\pi$ and $P^*P^*\pi$ generate the pion exchange potential which becomes important especially when P and P^* mesons are degenerate in the heavy quark limit. Here we note that only P meson cannot generate the pion exchange potential because the $PP\pi$ vertex is not allowed due to the parity invariance. In the literatures, this idea has been tested and shown to be indeed the case for heavy quark systems [10–13]. This does not seem to work, however, for light flavor sector in which the heavy quark symmetry is not a good symmetry [13].

In this paper, we study manifestly exotic baryons formed by a \bar{D} or B meson and a nucleon N , $\bar{D}N$ and BN , whose minimal quark content is $\bar{Q}qqqq$, where Q and q stand for heavy and light quarks, respectively [16]. In Refs. [12, 13], the investigation was made for the *negative* parity states where the bound and resonant states were discussed. Here we perform the analyses for the *positive* parity states to complete the investigation.

TABLE I. Various coupled channels for a given quantum number J^P for the positive parity $P = +1$.

J^P	channels
$1/2^+$	$PN(^2P_{1/2})$ $P^*N(^2P_{1/2})$ $P^*N(^4P_{1/2})$
$3/2^+$	$PN(^2P_{3/2})$ $P^*N(^2P_{3/2})$ $P^*N(^4P_{3/2})$ $P^*N(^4F_{3/2})$
$5/2^+$	$PN(^2F_{5/2})$ $P^*N(^4P_{5/2})$ $P^*N(^2F_{5/2})$ $P^*N(^4F_{5/2})$

This paper is organized as follows. In section 2, we briefly summarize the interaction between a heavy meson \bar{D} or B and a nucleon N by following the prescription in Refs. [12, 13]. In section 3, we solve the Schrödinger equations numerically and search the bound and resonant states in several quantum numbers (I, J^P) . Unlike the negative parity states, bound states are not found in the positive parity states but only resonances are. The difference of the present results from the previous ones is discussed. In the final section, we summarize the present work and discuss some future directions.

II. INTERACTIONS

Let us consider two-body states of a heavy meson and a nucleon with positive parity. Those states can be classified by isospin I , total spin J and parity P . In the present study, we consider the states with $I = 0$ and 1, and $J^P = 1/2^+$, $3/2^+$ and $5/2^+$, as summarized in Table I. In these systems, a heavy pseudoscalar meson (P) and a heavy vector meson (P^*) are degenerate in the heavy quark limit. Therefore, each state may contain both P and P^* , leading to a problem with coupled channels; three channels for $J^P = 1/2^+$ and four channels for $J^P = 3/2^+$ and $5/2^+$.

To obtain the interactions for a heavy meson and a nucleon, we employ Lagrangians satisfying the heavy quark symmetry and chiral symmetry [17, 18]. They are well-known and given as

$$\mathcal{L}_{\pi HH} = ig_\pi \text{Tr} [H_b \gamma_\mu \gamma_5 A_{ba}^\mu \bar{H}_a] , \quad (1)$$

$$\mathcal{L}_{v HH} = -i\beta \text{Tr} [H_b v^\mu (\rho_\mu)_{ba} \bar{H}_a] + i\lambda \text{Tr} [H_b \sigma^{\mu\nu} F_{\mu\nu}(\rho)_{ba} \bar{H}_a] , \quad (2)$$

where the subscripts π and v ($= \rho$ and ω) are for the pion and vector mesons. We consider not only the pion exchange which is relevant at long distances, but also the vector meson

TABLE II. Masses and coupling constants of mesons in Ref. [13].

	m_α [MeV]	g_π	β	λ [GeV $^{-1}$]	$g_{\alpha NN}^2/4\pi$	κ
π	137.27	0.59	—	—	13.6	—
ρ	769.9	—	0.9	0.56	0.84	6.1
ω	781.94	—	0.9	0.56	20.0	0.0

TABLE III. Cut-off parameters of a nucleon and heavy mesons in Ref. [13].

Potential	Λ_N [MeV]	Λ_D [MeV]	Λ_B [MeV]
π	830	1121	1070
$\pi\rho\omega$	846	1142	1091

exchange which is relevant at short distances. In Eq. (2), v^μ is the four-velocity of a heavy quark. In Eqs. (1) and (2), the heavy meson fields of $\bar{Q}q$ are parametrized by the heavy pseudoscalar and vector mesons,

$$H_a = \frac{1 + \not{v}}{2} [P_{a\mu}^* \gamma^\mu - P_a \gamma_5] , \quad (3)$$

$$\bar{H}_a = \gamma_0 H_a^\dagger \gamma_0 , \quad (4)$$

where the subscripts a, b are for light flavors, u, d . From Eqs. (1) and (2), we obtain the pion and vector meson vertices in the static approximation $v^\mu = (1, \vec{0})$. The coupling constants g_π, β and λ are the same as in our previous papers Refs. [12, 13] as summarized in Table II.

The interaction Lagrangians for a meson and nucleons are given by the standard form,

$$\mathcal{L}_{\pi NN} = \sqrt{2} i g_{\pi NN} \bar{N} \gamma_5 \hat{\pi} N , \quad (5)$$

$$\mathcal{L}_{v NN} = \sqrt{2} g_{v NN} \left[\bar{N} \gamma_\mu \hat{\rho}^\mu N + \frac{\kappa}{2m_N} \bar{N} \sigma_{\mu\nu} \partial^\nu \hat{\rho}^\mu N \right] , \quad (6)$$

where $N = (p, n)^T$ is the nucleon field. The coupling constants for the nucleon are taken from the phenomenological nuclear potential in Ref. [19] as summarized in Table II.

The potentials are derived by the vertices (1), (2), (5) and (6) as shown in Appendix. To take into account the internal structure of the hadrons, form factors associated with finite size of the mesons and nucleons are introduced at each vertex. We introduce the monopole type form factors as shown in Appendix. Here we have two cut-off parameters for \bar{D} (B)

mesons and a nucleon. The cut-off parameter for the nucleon is determined to reproduce the properties of the deuteron: the binding energy, scattering length and effective range. The cut-off parameters for \bar{D} (B) mesons are determined from the ratios of matter radii of \bar{D} (B) meson and nucleon, which are estimated by a quark model, as discussed in Ref. [12, 13]. Using those potentials, we solve the coupled-channel Schrödinger equations numerically. We employ the two potentials; the π exchange potential and the $\pi\rho\omega$ potential to discuss the important role of the pion. The cut-off parameters of the nucleon vertices for each parameter set are summarized in Table III.

III. SCATTERING STATES AND RESONANCES

After solving the Schrödinger equations, we find no bound state neither for $\bar{D}N$ nor for BN systems. However, by analyzing the scattering states, we find several resonances in the isosinglet channel ($I = 0$) both for $\bar{D}N$ and for BN , but no structure in the isotriplet channel ($I = 1$).

Let us first explain how resonant states are determined in the present study. The resonance energy and decay width are obtained by analyzing the phase shift in the scattering states. In the previous work [13], we identified a resonance at the position where the phase shift crosses $\pi/2$, because the decay width was relatively smaller than the resonance energy from the threshold. In the present study, because it will turn out that the decay widths are not necessarily small, we define the resonance position E_{re} by an inflection point of the phase shift [20]. Then, the width is obtained by $\Gamma = 2/(\partial\delta/\partial E)_{E=E_{\text{re}}}$ at the inflection point.

Now we discuss the result of each channel $(I, J^P) = (0, 1/2^+)$, $(0, 3/2^+)$ and $(0, 5/2^+)$. In the $(I, J^P) = (0, 1/2^+)$ channel, we find resonances in both $\bar{D}N$ and BN systems. The phase shifts δ 's of $PN(^2P_{1/2})$, $P^*N(^2P_{1/2})$ and $P^*N(^4P_{1/2})$ channels obtained for the $\pi\rho\omega$ potential are shown as functions of the scattering energy E in the center of mass system in Fig. 1. The vertical dashed lines in the figures represent the positions of \bar{D}^*N and B^*N thresholds. The resonance energies are measured from the lowest thresholds ($\bar{D}N$ and BN). The sharp increase of the phase shift of the $PN(^2P_{1/2})$ channel indicates the existence of a resonance. A similar behavior is obtained also when the π exchange potential is employed. The resonance energies and decay widths are summarized in Table IV. We obtain the resonance energy at 26.8 MeV for $\bar{D}N$ and at 5.8 MeV for BN , with the decay widths 131.3

MeV and 6.0 MeV, respectively, for the $\pi\rho\omega$ potential. Here, we compare two results by the π exchange potential and the $\pi\rho\omega$ potential, and find the difference is very small. As a result, the vector meson (ρ and ω) exchange interaction plays a minor role, while the π exchange interaction plays a dominant role to generate resonant states. In particular, the tensor force is important because no bound or resonant state exists without the PN - P^*N mixing. The pion dominance was also seen for the negative parity states as in our previous work [13]. The resonances are generated below the P^*N threshold, where the PN channel is open and the P^*N channel is closed. Here, we note that the PN - PN channel has no interaction in the pion exchange potential due to the parity non-conservation at the $PP\pi$ vertex, as presented explicitly in the π exchange potential (A3) in Appendix. Therefore the attractive force which forms the resonance in the PN channel is mainly provided from the PN - P^*N mixing effect. As a consequence, the mixing effect yields sufficient attraction to form the so-called shape resonance in the $PN(^2P_{1/2})$ channel with the p-wave centrifugal barrier. The total cross sections for the $\bar{D}N$ and BN scatterings when the $\pi\rho\omega$ potential is used are shown in Figs. 2 and 3, respectively. The peaks are found at around each resonance energy, 26.8 MeV and 5.8 MeV, for $\bar{D}N$ and BN , respectively.

TABLE IV. The resonance energies E_{re} and decay widths Γ for $(I, J^P) = (0, 1/2^+)$.

	$\bar{D}N(\pi)$	$\bar{D}N(\pi\rho\omega)$	$BN(\pi)$	$BN(\pi\rho\omega)$
E_{re} [MeV]	26.1	26.8	5.8	5.8
Γ [MeV]	125.2	131.3	5.8	6.0

In the $(I, J^P) = (0, 3/2^+)$ channel, we find a resonance for each $\bar{D}N$ and BN state. For $\bar{D}N$ state, we show the phase shifts δ 's of $\bar{D}N(^2P_{3/2})$, $\bar{D}^*N(^2P_{3/2})$, $\bar{D}^*N(^4P_{3/2})$ and $\bar{D}^*N(^4F_{3/2})$ in Fig. 4. There is a small peak structure in the phase shift of $\bar{D}N(^2P_{3/2})$ at the \bar{D}^*N threshold, which is interpreted as a cusp. On the other hand, the phase shift of $\bar{D}^*N(^4P_{3/2})$ which rises sharply indicates the presence of a resonance in this channel. Therefore the resonant state exists in the $\bar{D}^*N(^4P_{3/2})$ channel. The resonance energies and decay widths are summarized in Table V. When the PN - P^*N mixing is ignored, there still exists a resonance at the resonance energy 145.5 MeV with the decay width 6.1 MeV, which are close to the original values in the full channel-couplings. Therefore, the obtained resonance is a shape resonance generated mainly by the p-wave centrifugal barrier in the

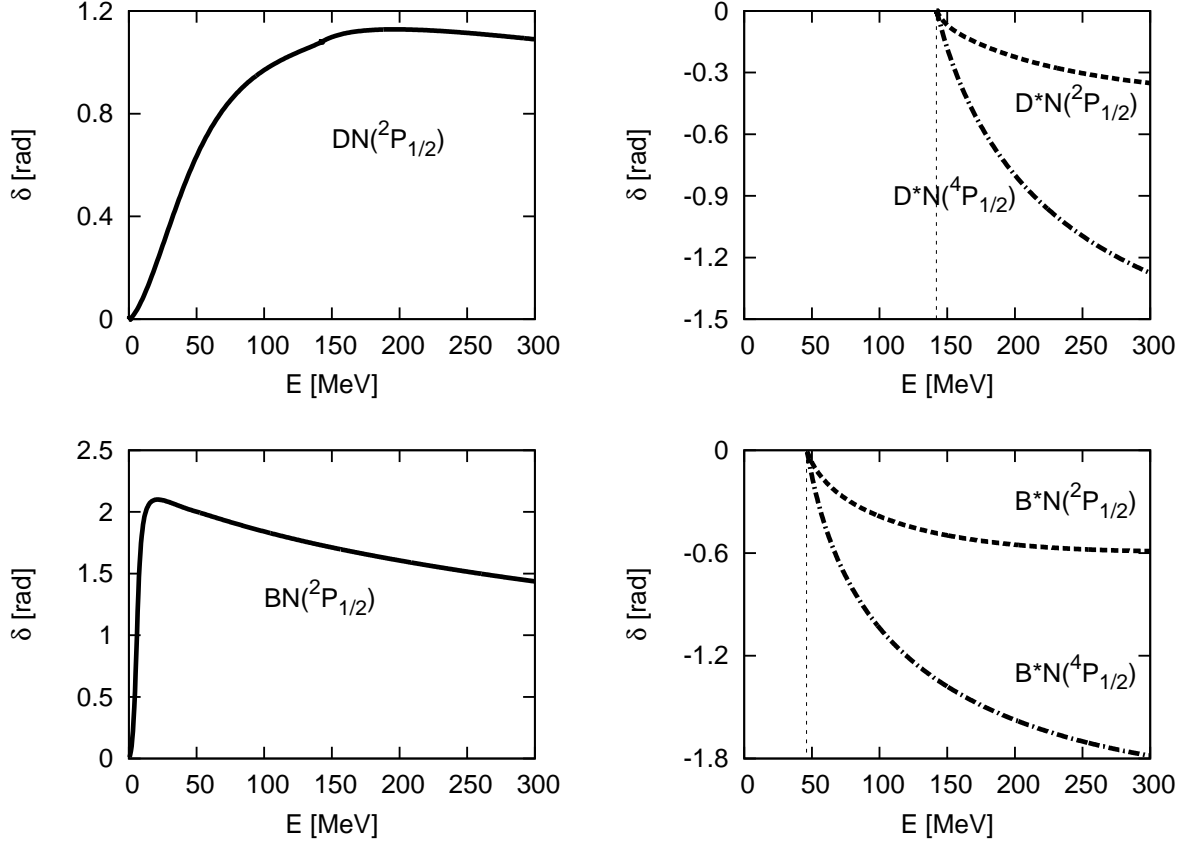


FIG. 1. Phase shifts of the $\bar{D}N$ and BN scattering states with $(I, J^P) = (0, 1/2^+)$ when the $\pi\rho\omega$ potential is used.

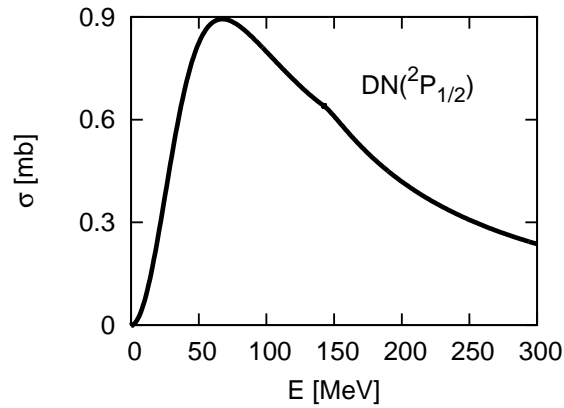


FIG. 2. Total cross section of the $\bar{D}N$ scattering state with $(I, J^P) = (0, 1/2^+)$ when the $\pi\rho\omega$ potential is used.

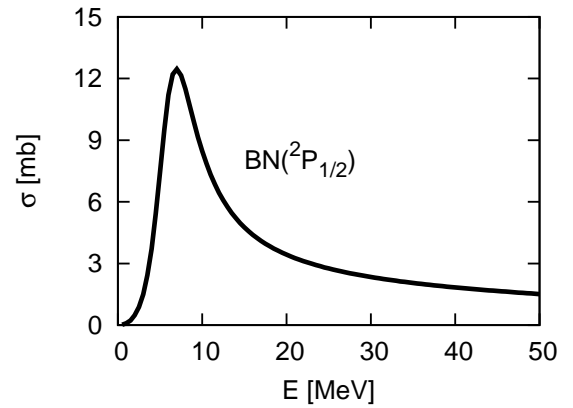


FIG. 3. Total cross section of the BN scattering state with $(I, J^P) = (0, 1/2^+)$ when the $\pi\rho\omega$ potential is used.

$\bar{D}^*N(^4P_{3/2})$ states in the P^*N channel. We show the cross section for $\bar{D}N(^4P_{3/2})$ for the $\pi\rho\omega$ potential in Fig. 6, in which the peak appears at the resonance energy 148.2 MeV.

For the BN state with $(I, J^P) = (0, 3/2^+)$, the phase shifts δ 's of $BN(^2P_{3/2})$, $B^*N(^2P_{3/2})$, $B^*N(^4P_{3/2})$ and $B^*N(^4F_{3/2})$ are plotted in Fig. 5. We find that the sharp increase of the phase shift passing through $\pi/2$ in $BN(^2P_{3/2})$ as an indication of a resonance. We also find that the phase shift in $B^*N(^4P_{3/2})$ starts from π . Therefore, the obtained resonance can be regarded as a bound state of B^*N . Indeed, when we switch off the $BN(^2P_{3/2})$ channel, we obtain a bound state of B^*N which energy is close to the original resonance position. Therefore, we conclude that the resonance in the BN state with $(I, J^P) = (0, 3/2^+)$ is a Feshbach resonance. The resonance energy is 31.8 MeV and the decay width is 28.7 MeV as summarized in Table V. In Fig. 7, we plot the cross section for BN with $(I, J^P) = (0, 3/2^+)$, where we see a peak at around the resonance energy 31.8 MeV.

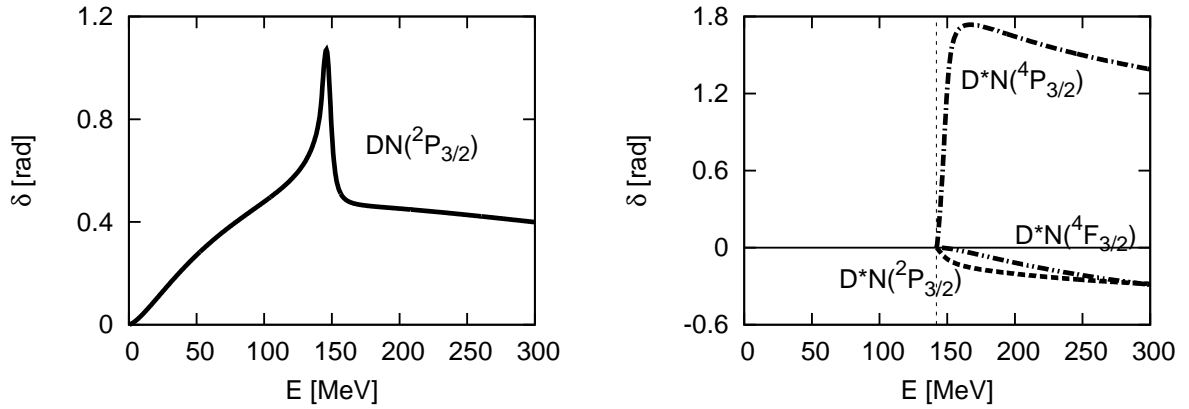


FIG. 4. Phase shifts of the $\bar{D}N$ scattering states with $(I, J^P) = (0, 3/2^+)$ when the $\pi\rho\omega$ potential is used.

TABLE V. The resonance energies E_{re} and decay widths Γ for $(I, J^P) = (0, 3/2^+)$.

	$\bar{D}^*N(^4P_{3/2})(\pi)$	$\bar{D}^*N(^4P_{3/2})(\pi\rho\omega)$	$BN(\pi)$	$BN(\pi\rho\omega)$
E_{re} [MeV]	148.2	148.2	32.3	31.8
Γ [MeV]	10.0	10.1	28.9	28.7

In the $(I, J^P) = (0, 5/2^+)$ channel, we find a resonance for each $\bar{D}N$ and BN above the P^*N threshold. The phase shifts δ 's of $PN(^2F_{5/2})$, $P^*N(^4F_{5/2})$, $P^*N(^2F_{5/2})$ and $P^*N(^4F_{5/2})$

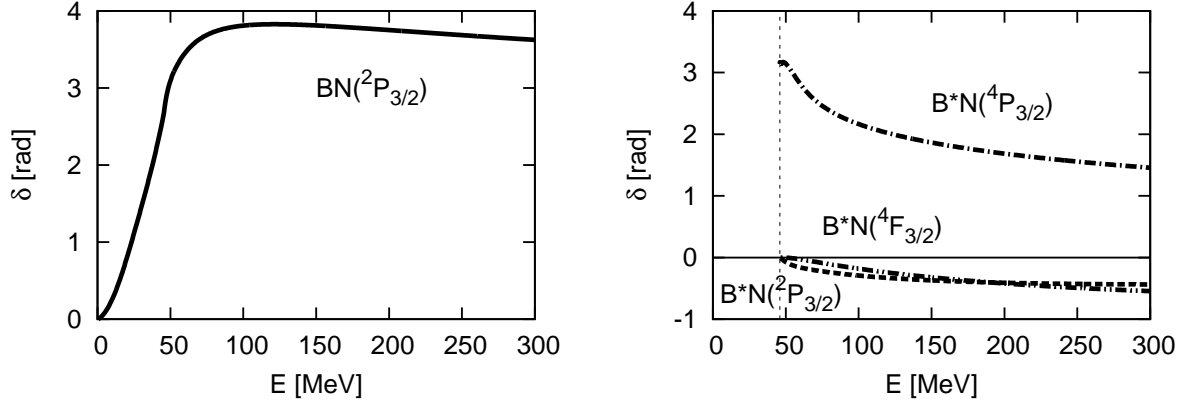


FIG. 5. Phase shifts of the BN scattering states with $(I, J^P) = (0, 3/2^+)$ when the $\pi\rho\omega$ potential is used.

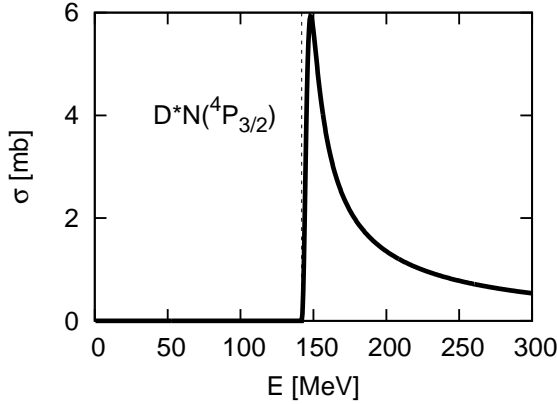


FIG. 6. Total cross section of the $\bar{D}^*N(^4P_{3/2})$ scattering when the $\pi\rho\omega$ potential is used.

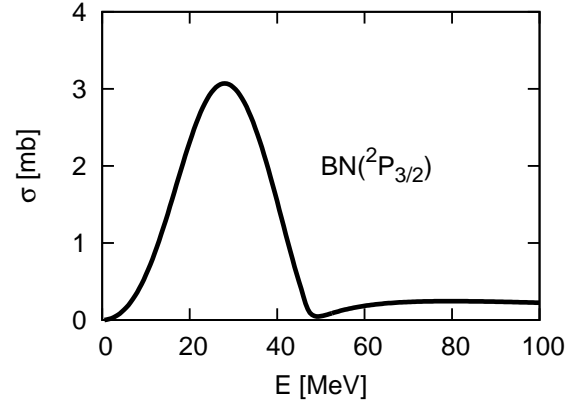


FIG. 7. Total cross section of the BN scattering state with $(I, J^P) = (0, 3/2^+)$ when the $\pi\rho\omega$ potential is used.

are shown in Fig. 8. Small peak structures in the phase shifts of $\bar{D}N(^2F_{5/2})$ and $BN(^2F_{5/2})$ are interpreted as cusps. Above the P^*N threshold, the phase shifts of $\bar{D}^*N(^4P_{5/2})$ and $B^*N(^4P_{5/2})$ rise up and these structures indicate the presence of resonances. The resonance energies are 176.0 MeV for \bar{D}^*N and 58.4 MeV for B^*N , and the decay widths are 174.8 MeV and 49.6 MeV, respectively. We summarize the results in Table VI. When the PN - P^*N mixing is ignored, the resonant states in the $\bar{D}^*N(^4P_{5/2})$ and $B^*N(^4P_{5/2})$ channels still exist at the resonance energies close to the values from the full channel-couplings. Therefore, these

resonant states are shape resonances generated mainly by the p-wave centrifugal barrier in the $P^*N(^4P_{5/2})$ channel. The cross sections for $\bar{D}^*N(^4P_{5/2})$ and $B^*N(^4P_{5/2})$ are plotted in Figs. 9 and 10, respectively.

Several comments are in order. First, in each (I, J^P) state, we verify that the results of the π exchange potential is similar to those of the $\pi\rho\omega$ potential. It means that both $\bar{D}N$ and BN systems are dominated almost by the long range force due to the pion exchange. Second, the mixing between PN and P^*N is important for the positive parity states, as for the negative parity states [12, 13]. The $\bar{D}N$ and BN resonances with $(I, J^P) = (0, 1/2^+)$ are generated with the p-wave centrifugal barrier in the PN channel by the attraction induced from the PN - P^*N mixing. The PN - P^*N mixing is important also for the BN resonant state with $(I, J^P) = (0, 3/2^+)$ because it is a Feshbach resonance. However, the $\bar{D}N$ resonance with $(I, J^P) = (0, 3/2^+)$ and the $\bar{D}N$ and BN resonances with $(I, J^P) = (0, 5/2^+)$ are generated with the p-wave centrifugal barrier in the P^*N channel mainly by the P^*N interaction, in which the PN - P^*N mixing effect plays a minor role. Therefore, in the positive parity states, resonances are generated by different mechanisms as summarized in Table VII. We note that the Feshbach resonance in the negative parity states was obtained both for $\bar{D}N$ and BN states with $(I, J^P) = (0, 3/2^-)$ in Ref. [13].

To compare the result of the positive parity states with the result of the negative parity states [12, 13], we show energy levels for the exotic states found in our investigations in Fig. 11. We find that the bound states exist in the negative parity states, while no bound state exists in the positive parity states. This is because the lowest state $PN(^2P_{1/2})$ in $J^P = 1/2^+$ has a p-wave orbital angular momentum $L = 1$, while $PN(^2S_{1/2})$ in $J^P = 1/2^-$ has an s-wave orbital angular momentum $L = 0$. In the same way, we see that the resonance energies of the positive parity states tend to be higher than those of the negative parity states.

TABLE VI. The resonance energies E_{re} and decay widths Γ for $(I, J^P) = (0, 5/2^+)$.

	$\bar{D}^*N(^4P_{5/2})(\pi)$	$\bar{D}^*N(^4P_{5/2})(\pi\rho\omega)$	$B^*N(^4P_{5/2})(\pi)$	$B^*N(^4P_{5/2})(\pi\rho\omega)$
E_{re} [MeV]	177.1	176.0	58.5	58.4
Γ [MeV]	184.6	174.8	52.2	49.6

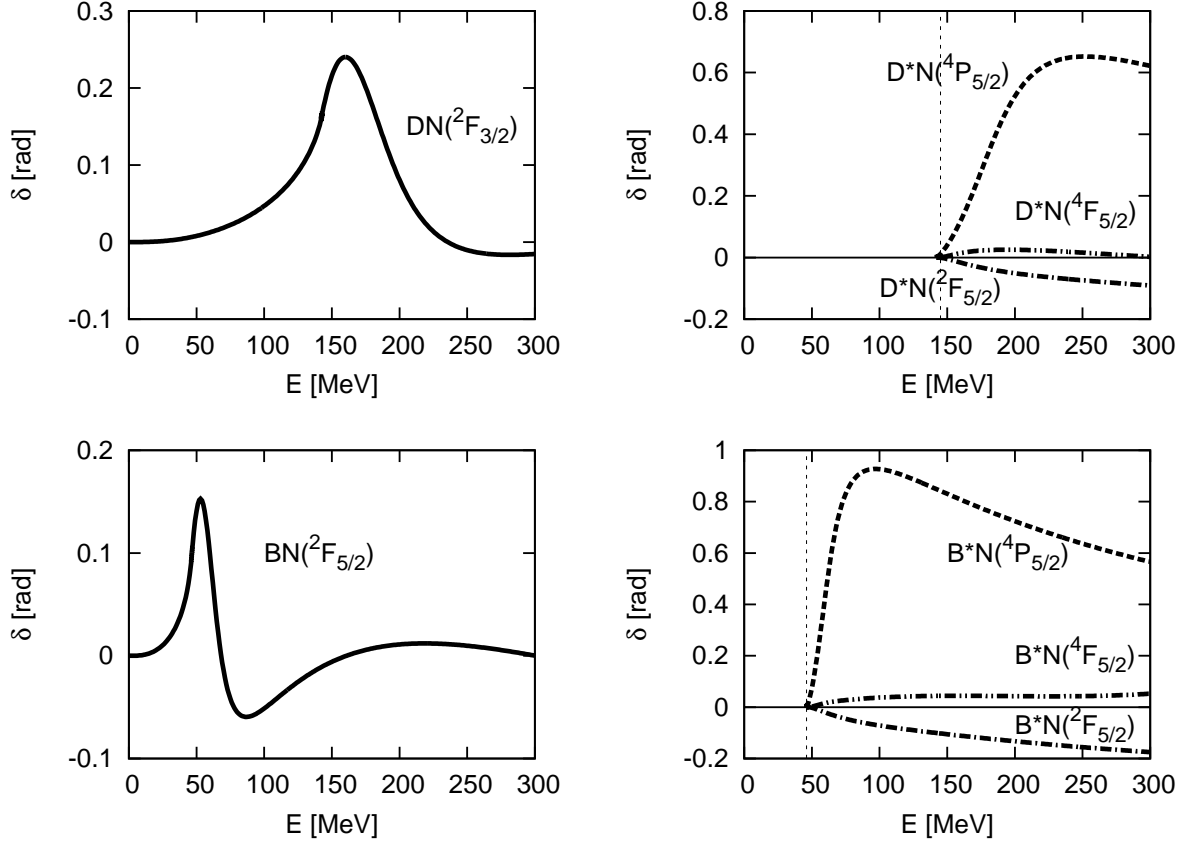


FIG. 8. Phase shifts of the $\bar{D}N$ and BN scattering states with $(I, J^P) = (0, 5/2^+)$ when the $\pi\rho\omega$ potential is used.

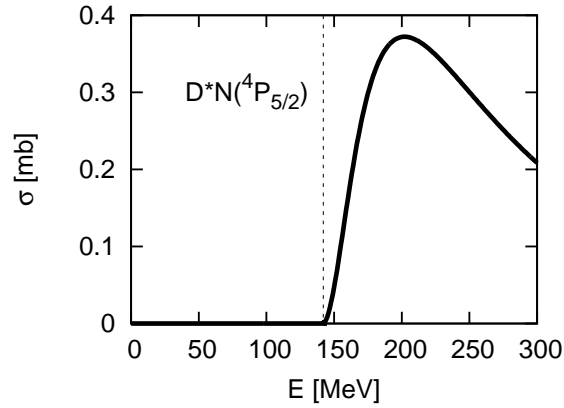


FIG. 9. Total cross section of the $\bar{D}^*N(^4P_{5/2})$ scattering when the $\pi\rho\omega$ potential is used.

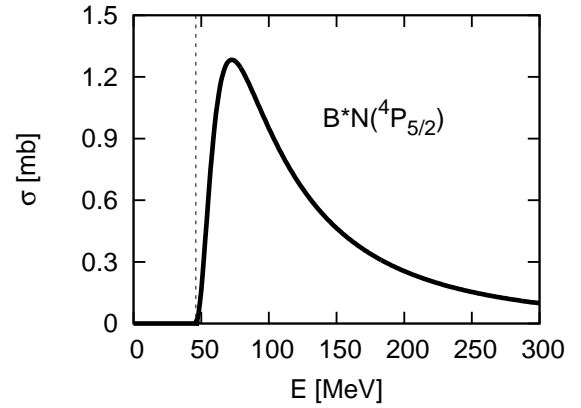


FIG. 10. Total cross section of the $B^*N(^4P_{5/2})$ scattering when the $\pi\rho\omega$ potential is used.

TABLE VII. The mechanism to form the resonances in the $\bar{D}N$ and BN states with $(I, J^P) = (0, 1/2^+)$, $(0, 3/2^+)$ and $(0, 5/2^+)$. All the shape resonances are induced by the p-wave centrifugal barrier in the PN and P^*N channels, respectively.

(I, J^P)	$\bar{D}N$ states	BN states
$(0, 1/2^+)$	shape resonance in PN	
$(0, 3/2^+)$	shape resonance in P^*N	Feshbach resonance
$(0, 5/2^+)$	shape resonance in P^*N	

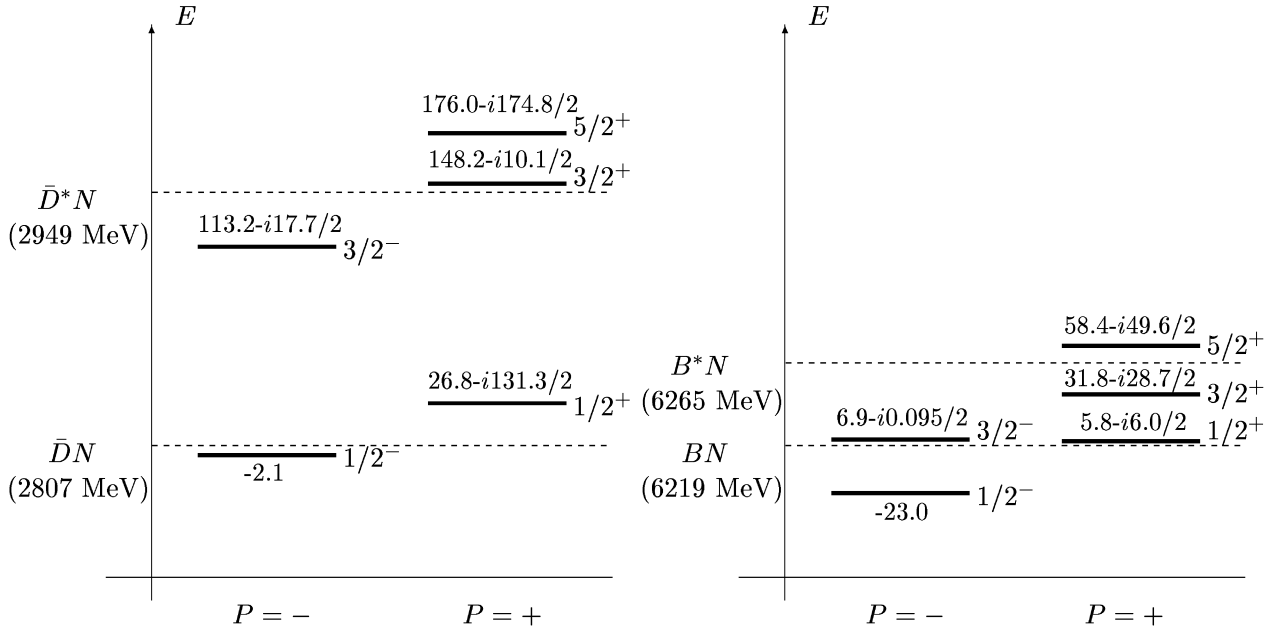


FIG. 11. Exotic states with positive parity ($P = +$) and negative parity ($P = -$). The energies are measured from the lowest thresholds ($\bar{D}N$ and BN). The binding energy is given as a real negative value, and the resonance energy E_{re} and decay width Γ are given as $E_{\text{re}} - i\Gamma/2$, in units of MeV. The values are given when the $\pi\rho\omega$ potential is used.

IV. SUMMARY

We have investigated exotic baryons constructed by a heavy meson and a nucleon like $\bar{D}N$ and BN . The interaction is given by the meson exchange potential between a $P = \bar{D}, B$ meson and a nucleon N , with respecting the heavy quark symmetry. To form those resonances, the PN - P^*N mixing originating from the heavy quark symmetry plays a crucial role. We have used the π exchange potential and the $\pi\rho\omega$ potential, and found that the pion exchange interaction works dominantly, while the vector meson exchange interaction plays only a minor role. Unlike the previous study for the negative parity case [12, 13], there is no bound state in the positive parity case. However, we have found new resonances for $(I, J^P) = (0, 1/2^+)$, $(0, 3/2^+)$ and $(0, 5/2^+)$ in isosinglet states, and no structure in isotriplet states. These resonant states exist near or below the thresholds of P^*N . These exotic systems will be interesting objects which can be searched and studied at accelerator facilities, such as J-PARC, FAIR and so on, and also in the relativistic heavy ion collisions at RHIC and LHC [21, 22].

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Appendix A: Potentials and kinetic terms

The interaction potentials are derived by using the Lagrangians Eqs. (1)-(6). In deriving the potentials we use the static approximation where the energy transfer can be neglected as compared to the momentum transfer. The resulting potentials for the coupled channel systems are given in the matrix form of 3×3 for $J^P = 1/2^+$ and of 4×4 for $J^P = 3/2^+$ and $5/2^+$,

$$V_{1/2^+} = \begin{pmatrix} V_{1/2^+}^{11} & V_{1/2^+}^{12} & V_{1/2^+}^{13} \\ V_{1/2^+}^{21} & V_{1/2^+}^{22} & V_{1/2^+}^{23} \\ V_{1/2^+}^{31} & V_{1/2^+}^{32} & V_{1/2^+}^{33} \end{pmatrix}, \quad (\text{A1})$$

$$V_{3/2^+, 5/2^+} = \begin{pmatrix} V_{3/2^+, 5/2^+}^{11} & V_{3/2^+, 5/2^+}^{12} & V_{3/2^+, 5/2^+}^{13} & V_{3/2^+, 5/2^+}^{14} \\ V_{3/2^+, 5/2^+}^{21} & V_{3/2^+, 5/2^+}^{22} & V_{3/2^+, 5/2^+}^{23} & V_{3/2^+, 5/2^+}^{24} \\ V_{3/2^+, 5/2^+}^{31} & V_{3/2^+, 5/2^+}^{32} & V_{3/2^+, 5/2^+}^{33} & V_{3/2^+, 5/2^+}^{34} \\ V_{3/2^+, 5/2^+}^{41} & V_{3/2^+, 5/2^+}^{42} & V_{3/2^+, 5/2^+}^{43} & V_{3/2^+, 5/2^+}^{44} \end{pmatrix}, \quad (\text{A2})$$

in the basis given in Table I in the same ordering. The π exchange potential between a heavy meson and a nucleon is obtained by

$$V_{1/2^+}^\pi = \frac{g_\pi g_{\pi NN}}{\sqrt{2} m_N f_\pi} \frac{1}{3} \begin{pmatrix} 0 & \sqrt{3} C_{m_\pi} & -\sqrt{6} T_{m_\pi} \\ \sqrt{3} C_{m_\pi} & -2 C_{m_\pi} & -\sqrt{2} T_{m_\pi} \\ -\sqrt{6} T_{m_\pi} & -\sqrt{2} T_{m_\pi} & C_{m_\pi} - 2 T_{m_\pi} \end{pmatrix} \vec{\tau}_P \cdot \vec{\tau}_N, \quad (\text{A3})$$

$$V_{3/2^+}^\pi = \frac{g_\pi g_{\pi NN}}{\sqrt{2} m_N f_\pi} \frac{1}{3} \begin{pmatrix} 0 & \sqrt{3} C_{m_\pi} & \sqrt{\frac{3}{5}} T_{m_\pi} & -3 \sqrt{\frac{3}{5}} T_{m_\pi} \\ \sqrt{3} C_{m_\pi} & -2 C_{m_\pi} & \frac{1}{\sqrt{5}} T_{m_\pi} & -\frac{3}{\sqrt{5}} T_{m_\pi} \\ \sqrt{\frac{3}{5}} T_{m_\pi} & \frac{1}{\sqrt{5}} T_{m_\pi} & C_{m_\pi} + \frac{8}{5} T_{m_\pi} & \frac{6}{5} T_{m_\pi} \\ -3 \sqrt{\frac{3}{5}} T_{m_\pi} & -\frac{3}{\sqrt{5}} T_{m_\pi} & \frac{6}{5} T_{m_\pi} & C_{m_\pi} - \frac{8}{5} T_{m_\pi} \end{pmatrix} \vec{\tau}_P \cdot \vec{\tau}_N, \quad (\text{A4})$$

$$V_{5/2^+}^\pi = \frac{g_\pi g_{\pi NN}}{\sqrt{2}m_N f_\pi} \frac{1}{3} \begin{pmatrix} 0 & \frac{3}{5}\sqrt{10}T_{m_\pi} & \sqrt{3}C_{m_\pi} & -2\sqrt{\frac{3}{5}}T_{m_\pi} \\ \frac{3}{5}\sqrt{10}T_{m_\pi} & C_{m_\pi} - \frac{2}{5}T_{m_\pi} & \sqrt{\frac{6}{5}}T_{m_\pi} & \frac{4}{5}\sqrt{6}T_{m_\pi} \\ \sqrt{3}C_{m_\pi} & \sqrt{\frac{6}{5}}T_{m_\pi} & -2C_{m_\pi} & -\frac{2}{\sqrt{5}}T_{m_\pi} \\ -2\sqrt{\frac{3}{5}}T_{m_\pi} & \frac{4}{5}\sqrt{6}T_{m_\pi} & -\frac{2}{\sqrt{5}}T_{m_\pi} & C_{m_\pi} + \frac{2}{5}T_{m_\pi} \end{pmatrix} \vec{\tau}_P \cdot \vec{\tau}_N, \quad (\text{A5})$$

where $C_m = C(r; m)$, $T_m = T(r; m)$, and $\vec{\tau}_P$ and $\vec{\tau}_N$ are the isospin matrices for $P(P^*)$ and N . The functions $C(r; m)$ and $T(r; m)$ are given by

$$C(r; m) = \int \frac{d^3q}{(2\pi)^3} \frac{m^2}{\vec{q}^2 + m^2} e^{i\vec{q}\cdot\vec{r}} F(\Lambda_P, \vec{q}) F(\Lambda_N, \vec{q}), \quad (\text{A6})$$

$$T(r; m) S_{12}(\hat{r}) = \int \frac{d^3q}{(2\pi)^3} \frac{-\vec{q}^2}{\vec{q}^2 + m^2} S_{12}(\hat{q}) e^{i\vec{q}\cdot\vec{r}} F(\Lambda_P, \vec{q}) F(\Lambda_N, \vec{q}), \quad (\text{A7})$$

with $S_{12}(\hat{x}) = 3(\vec{\sigma}_1 \cdot \hat{x})(\vec{\sigma}_2 \cdot \hat{x}) - \vec{\sigma}_1 \cdot \vec{\sigma}_2$, and $F(\Lambda, \vec{q})$ denotes the form factor given by

$$F_\alpha(\Lambda, \vec{q}) = \frac{\Lambda^2 - m_\alpha^2}{\Lambda^2 + |\vec{q}|^2} \quad (\text{A8})$$

where m_α and \vec{q} are the mass and three-momentum of the incoming meson $\alpha (= \pi, \rho, \omega)$.

The corresponding potentials of the ρ meson exchange are given by

$$V_{1/2^+}^\rho = \frac{g_V g_{\rho NN} \beta}{\sqrt{2}m_\rho^2} \begin{pmatrix} C_{m_\rho} & 0 & 0 \\ 0 & C_{m_\rho} & 0 \\ 0 & 0 & C_{m_\rho} \end{pmatrix} \vec{\tau}_P \cdot \vec{\tau}_N + \frac{g_V g_{\rho NN} \lambda(1 + \kappa)}{\sqrt{2}m_N} \frac{1}{3} \begin{pmatrix} 0 & 2\sqrt{3}C_{m_\rho} & \sqrt{6}T_{m_\rho} \\ 2\sqrt{3}C_{m_\rho} & -4C_{m_\rho} & \sqrt{2}T_{m_\rho} \\ \sqrt{6}T_{m_\rho} & \sqrt{2}T_{m_\rho} & 2C_{m_\rho} + 2T_{m_\rho} \end{pmatrix} \vec{\tau}_P \cdot \vec{\tau}_N, \quad (\text{A9})$$

$$\begin{aligned}
V_{3/2^+}^\rho = & \frac{g_V g_{\rho NN} \beta}{\sqrt{2} m_\rho^2} \begin{pmatrix} C_{m_\rho} & 0 & 0 & 0 \\ 0 & C_{m_\rho} & 0 & 0 \\ 0 & 0 & C_{m_\rho} & 0 \\ 0 & 0 & 0 & C_{m_\rho} \end{pmatrix} \vec{\tau}_P \cdot \vec{\tau}_N \\
& + \frac{g_V g_{\rho NN} \lambda (1 + \kappa)}{\sqrt{2} m_N} \frac{1}{3} \begin{pmatrix} 0 & 2\sqrt{3}C_{m_\rho} & -\sqrt{\frac{3}{5}}T_{m_\rho} & 3\sqrt{\frac{3}{5}}T_{m_\rho} \\ 2\sqrt{3}C_{m_\rho} & -4C_{m_\rho} & -\frac{1}{\sqrt{5}}T_{m_\rho} & \frac{3}{\sqrt{5}}T_{m_\rho} \\ -\sqrt{\frac{3}{5}}T_{m_\rho} & -\frac{1}{\sqrt{5}}T_{m_\rho} & 2C_{m_\rho} - \frac{8}{5}T_{m_\rho} & -\frac{6}{5}T_{m_\rho} \\ 3\sqrt{\frac{3}{5}}T_{m_\rho} & \frac{3}{\sqrt{5}}T_{m_\rho} & -\frac{6}{5}T_{m_\rho} & 2C_{m_\rho} + \frac{8}{5}T_{m_\rho} \end{pmatrix} \vec{\tau}_P \cdot \vec{\tau}_N,
\end{aligned} \tag{A10}$$

$$\begin{aligned}
V_{5/2^+}^\rho = & \frac{g_V g_{\rho NN} \beta}{\sqrt{2} m_\rho^2} \begin{pmatrix} C_{m_\rho} & 0 & 0 & 0 \\ 0 & C_{m_\rho} & 0 & 0 \\ 0 & 0 & C_{m_\rho} & 0 \\ 0 & 0 & 0 & C_{m_\rho} \end{pmatrix} \vec{\tau}_P \cdot \vec{\tau}_N \\
& + \frac{g_V g_{\rho NN} \lambda (1 + \kappa)}{\sqrt{2} m_N} \frac{1}{3} \begin{pmatrix} 0 & -\frac{3}{5}\sqrt{10}T_{m_\rho} & 2\sqrt{3}C_{m_\rho} & 2\sqrt{\frac{3}{5}}T_{m_\rho} \\ -\frac{3}{5}\sqrt{10}T_{m_\rho} & 2C_{m_\rho} + \frac{2}{5}T_{m_\rho} & -\sqrt{\frac{6}{5}}T_{m_\rho} & -\frac{4}{5}\sqrt{6}T_{m_\rho} \\ 2\sqrt{3}C_{m_\rho} & -\sqrt{\frac{6}{5}}T_{m_\rho} & -4C_{m_\rho} & \frac{2}{\sqrt{5}}T_{m_\rho} \\ 2\sqrt{\frac{3}{5}}T_{m_\rho} & -\frac{4}{5}\sqrt{6}T_{m_\rho} & \frac{2}{\sqrt{5}}T_{m_\rho} & 2C_{m_\rho} - \frac{2}{5}T_{m_\rho} \end{pmatrix} \vec{\tau}_P \cdot \vec{\tau}_N.
\end{aligned} \tag{A11}$$

The ω meson exchange potential can be obtained by replacing the relevant coupling constants and the mass of the exchanged meson, and by removing the isospin factor $\vec{\tau}_P \cdot \vec{\tau}_N$. The anomalous coupling κ for the ω meson exchange potential is set as zero in Eqs. (A9)-(A11).

The kinetic terms are given by

$$K_{1/2^+} = \text{diag} \left(-\frac{1}{2\tilde{m}_P} \Delta_1, -\frac{1}{2\tilde{m}_{P^*}} \Delta_1 + \Delta m_{PP^*}, -\frac{1}{2\tilde{m}_{P^*}} \Delta_1 + \Delta m_{PP^*} \right), \quad (\text{A12})$$

$$K_{3/2^+} = \text{diag} \left(-\frac{1}{2\tilde{m}_P} \Delta_1, -\frac{1}{2\tilde{m}_{P^*}} \Delta_1 + \Delta m_{PP^*}, -\frac{1}{2\tilde{m}_{P^*}} \Delta_1 + \Delta m_{PP^*}, \right. \\ \left. -\frac{1}{2\tilde{m}_{P^*}} \Delta_3 + \Delta m_{PP^*} \right), \quad (\text{A13})$$

$$K_{5/2^+} = \text{diag} \left(-\frac{1}{2\tilde{m}_P} \Delta_3, -\frac{1}{2\tilde{m}_{P^*}} \Delta_1 + \Delta m_{PP^*}, -\frac{1}{2\tilde{m}_{P^*}} \Delta_3 + \Delta m_{PP^*}, \right. \\ \left. -\frac{1}{2\tilde{m}_{P^*}} \Delta_3 + \Delta m_{PP^*} \right), \quad (\text{A14})$$

for $J^P = 1/2^+, 3/2^+$ and $5/2^+$, respectively. Here, we define $\Delta_l = \partial^2/\partial r^2 + (2/r)\partial/\partial r - l(l+1)/r^2$ and $\tilde{m}_{P^{(*)}} = m_N m_{P^{(*)}}/(m_N + m_{P^{(*)}})$, with $\Delta m_{PP^*} = m_{P^*} - m_P$. The total Hamiltonian is then given by $H_{J^P} = K_{J^P} + V_{J^P}$.

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